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$B \to K(K^*)$ + missing energy in unparticle physics

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Abstract: In the present work we study the effects of an unparticle $\mathcal{U}$ as the possible source of missing energy in the decay $B \to K(K^*)$ + missing energy. We find that the dependence of the differential branching ratio on the $K(K^*)$-meson's energy in the presence of the vector unparticle operators is very distinctive from that of the SM. Moreover, in using the existing upper bound on $B \to K(K^*)$ + missing energy decays, we have been able to put more stringent constraints on the parameters of unparticle stuff.

Keywords: Beyond Standard Model, Rare Decays, B-Physics.

Contents

1. Introduction

Flavour Changing Neutral Current (FCNC) processes are not only powerful tests of the Standard Model (SM) but also provide very stringent tests for any physics beyond it. The smallness of FCNC processes in the SM is attributed to the fact that these processes are generated at loop level and are further suppressed by the CKM factors. Due to their smallness within the SM these processes can also be very sensitive to any new physics beyond the SM. Amongst the many FCNC decays involving B and K-mesons the decays of the form $b \to s +$ missing energy have been the focus of much investigation at the B factories Belle and Babar.

Of particular interest, in the SM, is the decay $b \to s\nu\bar{\nu}$, as it has the theoretical advantage of uncertainties much smaller than those of other decays, due to the absence of a photonic penguin contribution and hadronic long distance effects. However, in spite these theoretical advantages, it might be very difficult to measure the inclusive mode $B \to X_s\nu\bar{\nu}$, as it requires a construction of all the X_s's. Therefore the rare $B \to K(K^*)\nu\bar{\nu}$ decays play a special role, both from experimental and theoretical points of view. Also the branching fractions of the B-meson decays are quite large, with theoretical estimates of $Br(B \to K^*\nu\bar{\nu}) \sim 10^{-5}$ and $Br(B \to K\nu\bar{\nu}) \sim 10^{-6}$ [1]. These processes, based on $b \to s\nu\bar{\nu}$, are very sensitive to non-standard Z models and have been extensively studied in the literature [2–4].

As such, any new physics model which can provide a relatively light new source of *missing energy* can potentially enhance the observed rates of $B \to K(K^*)+$missing energy $(B \to K(K^*) + \not{E})$, where many models have been proposed which provide such low mass candidates (which can contribute to $b \to s +$ missing energy). Note that in reference [3] the phenomenology of such low mass scalars was explored. Such studies have also been done in the context of large extra dimension models [5] and leptophobic Z' models [1, 2]. One such model, which has excited much interest recently, is that of Unparticles, as proposed by H. Georgi [6]. In this model we assume that at a very high energy our theory contains both the fields of the SM and the fields of a theory with a nontrivial IR fixed point, which he called the Banks-Zaks (BZ) fields [7]. In his model these two sets interacted through the exchange of particles with a large mass scale $M_{\mathcal{U}}$, where below this scale there were nonrenormablizable couplings between the SM fields and the BZ fields suppressed by powers of $M_{\mathcal{U}}$. The renormalizable couplings of the BZ fields then produced dimensional transmutation, and the scale-invariant unparticle fields emerged below an energy scale $\Lambda_{\mathcal{U}}$.

In the effective theory below $\Lambda_{\mathcal{U}}$ the BZ operators matched onto the unparticle operators, and the nonrenormaliable interactions matched onto a new set of interactions between the SM and unparticle fields. The end result was a collection of unparticle stuff with scale dimension $d_{\mathcal{U}}$, which looked like a non-integral number $d_{\mathcal{U}}$ of invisible massless particles, whose production might be detectable in missing energy and momentum distributions [6].

Recently there has been a lot of interest in unparticle physics [6, 8−16], where the signatures of *unparticles* have been discussed at colliders [8, 10, 15], in Lepton Flavor Violating (LFV) processes [13], cosmology and astrophysics [16], and low energy processes [11, 12, 9].

In the present work we study the $B \to K(K^*) + \not{E}$ decay in unparticle theory, where this work is organized as follows: In section 2 we calculate the various contributions from both the SM and unparticle theory to the above-mentioned decays. Section 3 contains our numerical analysis and conclusions.

2. Differential decay widths

In the SM the decay mode $B \to K(K^*) + \not{E}$ is described by the decay $B \to K(K^*)\nu\bar{\nu}$. As was noted earlier, unparticles can also contribute to these decays. Hence a comparison of the signatures of the two decay modes $B \to K(K^*)\nu\bar{\nu}$ and $B \to K(K^*)\mathcal{U}$ is required.

In the SM the decay $B \to K(K^*)\nu\bar{\nu}$ is described by the quark level process $b \to s\nu\bar{\nu}$ through the effective Hamiltonian:

$$\mathcal{H} = \frac{G_F}{\sqrt{2}}\frac{\alpha}{2\pi}V_{tb}V_{ts}^{*}C_{10}\ \bar{s}\gamma_\mu\left(1-\gamma_5\right)b\ \bar{\nu}\gamma_\mu\left(1-\gamma_5\right)\nu\ , \tag{2.1}$$

where

$$C_{10} = \frac{X(x_t)}{sin^2\theta_w}\ , \tag{2.2}$$

and the $X(x_t)$ is the usual Inami-Lim function, given as:

$$X(x_t) = \frac{x_t}{8}\left\{\frac{x_t+1}{x_t-1} + \frac{3x_t-6}{(x_t-1)^2}ln(x_t)\right\}\ , \tag{2.3}$$

with $x_t = m_t/m_W^2$.

Similarly, the unparticle transition at quark level can be described by $b \to s\mathcal{U}$, where we shall consider the following operators:

$$\begin{aligned} \text{Scalar unparticle operators} &\implies \mathcal{C}_S \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}\ \bar{s}\gamma_\mu b\ \partial^\mu O_{\mathcal{U}} + \mathcal{C}_P \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}\ \bar{s}\gamma_\mu\gamma_5 b\ \partial^\mu O_{\mathcal{U}}\ , \\ \text{Vector unparticle operators} &\implies \mathcal{C}_V \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}}\ \bar{s}\gamma_\mu b\ O_{\mathcal{U}}^\mu\ +\ \mathcal{C}_A \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}}\ \bar{s}\gamma_\mu\gamma_5 b\ O_{\mathcal{U}}^\mu . \end{aligned} \tag{2.4}$$

Before proceeding with our analysis note that we shall write the propagator for the scalar unparticle field as [8, 10]:

$$\int d^4x e^{iP.x}\langle 0|TO_{\mathcal{U}}(x)O_{\mathcal{U}}(0)|0\rangle = i\frac{A_{d_{\mathcal{U}}}}{2\,\sin(d_{\mathcal{U}}\pi)}(-P^2)^{d_{\mathcal{U}}-2}\ , \tag{2.5}$$

where

$$A_{d_{\mathcal{U}}} = \frac{16\pi^{5/2}}{(2\pi)^{2d_{\mathcal{U}}}}\frac{\Gamma(d_{\mathcal{U}}+1/2)}{\Gamma(d_{\mathcal{U}}-1)\Gamma(2d_{\mathcal{U}})}\ .$$

2.1 The Standard Model

Using the SM effective Hamiltonian for the quark level process $b \to s\nu\bar{\nu}$, as given in equation (2.1), we can calculate the differential decay width of $B \to K(K^*)\nu\bar{\nu}$ (using the form factor definitions for the $B \to K$ transition as given in appendix A.1).

After taking into account the three species of SM neutrinos, we evaluate the differential decay width as a function of K-meson energy (E_K) as:

$$\frac{d\Gamma^{\rm SM}}{dE_K} = \frac{G_F^2\alpha^2}{2^7\pi^5 m_B^2}\ |V_{ts}V_{tb}^*|^2\ |C_{10}|^2 f_+^2(q^2)\ \lambda^{3/2}(m_B^2, m_K^2, q^2)\ , \tag{2.6}$$

where $\lambda(m_B^2, m_K^2, q^2) = m_B^4 + m_K^4 + q^4 - 2m_B^2 q^2 - 2m_K^2 q^2 - 2m_K^2 m_B^2$, and $q^2 = m_B^2 + m_K^2 - 2m_B E_K$.

Similarly, for the $B \to K^*$ case, using the definition of form factors for $B \to K^*$ transitions as given in appendix A.2, the differential decay rate in the SM can be calculated as:

$$\begin{aligned} \frac{d\Gamma^{\rm SM}}{dE_{K^*}} = \frac{G_F^2\alpha^2}{2^9\pi^5 m_B^2}|V_{ts}V_{tb}^*|^2\lambda^{1/2}|C_{10}|^2\Bigg(8\lambda q^2\frac{V^2}{(m_B+m_{K^*})^2} + \frac{1}{m_{K^*}^2}\Bigg[\lambda^2\frac{A_2^2}{(m_B+m_{K^*})^2} \\ +(m_B+m_{K^*})^2(\lambda+12m_{K^*}^2q^2)A_1^2 - 2\lambda(m_B^2-m_{K^*}^2-q^2)Re(A_1^*A_2)\Bigg]\Bigg)\ , \end{aligned} \tag{2.7}$$

where $\lambda = m_B^4 + m_{K^*}^4 + q^4 - 2m_B^2 q^2 - 2m_{K^*}^2 q^2 - 2m_{K^*}^2 m_B^2$, and $q^2 = m_B^2 + m_{K^*}^2 - 2m_B E_{K^*}$.

2.2 The scalar unparticle operator

As listed earlier, the following scalar operators can contribute to the $B \to K(K^*)\,\mathcal{U}$ decay:

$$\mathcal{C}_S\frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}\ \bar{s}\gamma_\mu b\ \partial^\mu O_{\mathcal{U}} + \mathcal{C}_P\frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}\ \bar{s}\gamma_\mu\gamma_5 b\ \partial^\mu O_{\mathcal{U}} = \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}\ \bar{s}\gamma_\mu\,(\mathcal{C}_S+\mathcal{C}_P\gamma_5)\,b\ \partial^\mu O_{\mathcal{U}}\ , \tag{2.8}$$

where we have defined our form factors in appendix A. As such, the matrix element for the process $B(p) \to K(p') + \mathcal{U}(q)$ can be written as:

$$\mathcal{M}^S = \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \, \mathcal{C}_S \left[f_+(m_B^2 - m_K^2) + f_- q^2 \right] O_{\mathcal{U}} \ . \tag{2.9}$$

The decay rate for $B(p) \to K(p')\mathcal{U}(q)$ can now be evaluated to be:

$$\frac{d\Gamma^{S\mathcal{U}}}{dE_K} = \frac{1}{8\pi^2 m_B} \frac{A_{d_{\mathcal{U}}}}{\Lambda_{\mathcal{U}}^{2d_{\mathcal{U}}}} \, |\mathcal{C}_S|^2 \sqrt{E_K^2 - m_K^2} \left(m_B^2 + m_K^2 - 2m_B E_K\right)^{d_{\mathcal{U}}-2} \times \left[f_+(m_B^2 - m_K^2) + f_-(m_B^2 + 2m_K^2 - 2m_B E_K) \right]^2 \ . \tag{2.10}$$

For the $B \to K^*$ transition our calculation proceeds along the same lines, where the matrix element for $B(p) \to K^*(p')\mathcal{U}(q)$ can be written as:

$$\mathcal{M}^S = \frac{i\mathcal{C}_P}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}(\epsilon.q) \left\{ (m_B + m_{K^*})A_1 - (m_B - m_{K^*})A_2 - 2m_{K^*}(A_3 - A_0) \right\} \ O_{\mathcal{U}} \ , \tag{2.11}$$

and the differential decay rate as:

$$\frac{d\Gamma^{S\mathcal{U}}}{dE_{K^*}} = \frac{m_B}{2\pi^2} \frac{A_{d_{\mathcal{U}}}}{\Lambda_{\mathcal{U}}^{2d_{\mathcal{U}}}} \, |\mathcal{C}_P|^2 A_0^2 \, \left(E_{K^*}^2 - m_{K^*}^2\right)^{3/2} \left(m_B^2 + m_{K^*}^2 - 2m_B E_{K^*}\right)^{d_{\mathcal{U}}-2} \ . \tag{2.12}$$

As can seen from the above expressions the scalar unparticle contribution to the decay rate for $B \to K\mathcal{U}$ and $B \to K^*\mathcal{U}$ will depend upon $\mathcal{C}_S$ and $\mathcal{C}_P$ respectively. This shall allow us to place constraints upon $\mathcal{C}_S$ and $\mathcal{C}_P$ from these two different decay modes. This issue shall be re-visited in the final section of this paper.

2.3 The vector unparticle operator

Along similar lines as followed in the previous subsection, we shall now make use of the vector unparticle operators:

$$\mathcal{C}_V \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \, \bar{s}\gamma_\mu b \, O_{\mathcal{U}}^\mu \ + \ \mathcal{C}_A \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \, \bar{s}\gamma_\mu\gamma_5 b \, O_{\mathcal{U}}^\mu = \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \, \bar{s}\gamma_\mu \left(\mathcal{C}_V + \mathcal{C}_A\gamma_5\right) b \, O_{\mathcal{U}}^\mu \ ,$$

and the form factors of appendix A, to calculate the matrix element for $B(p) \to K(p')\mathcal{U}(q)$:

$$\mathcal{M}^V = \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \, \mathcal{C}_V \left[f_+(p + p')_\mu + f_-(p - p')_\mu \right] O_{\mathcal{U}}^\mu \ . \tag{2.13}$$

And as such we calculate the differential decay rate as:

$$\frac{d\Gamma^{V\mathcal{U}}}{dE_K} = \frac{1}{8\pi^2 m_B} \frac{A_{d_{\mathcal{U}}}}{\Lambda_{\mathcal{U}}^{2d_{\mathcal{U}}-2}} \, |\mathcal{C}_V|^2 |f_+|^2 \left(m_B^2 + m_K^2 - 2m_B E_K\right)^{d_{\mathcal{U}}-2} \sqrt{E_K^2 - m_K^2} \times \left\{ -(m_B^2 + m_K^2 + 2m_B E_K) + \frac{(m_B^2 - m_K^2)^2}{(m_B^2 + m_K^2 - 2m_B E_K)} \right\}. \tag{2.14}$$

For the $B \to K^*$ case the matrix element for $B(p) \to K^*(p')\mathcal{U}(q)$ is:

$$\mathcal{M}^V = \left\{ \frac{\mathcal{C}_A}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \left(i\epsilon_\mu (m_B+m_{K^*}) A_1 - i(p+p')_\mu (\epsilon.q) \frac{A_2}{m_B+m_{K^*}} - iq_\mu (\epsilon.q) \frac{2m_{K^*}}{q^2} [A_3 - A_0] \right) \right. \\ \left. + \frac{\mathcal{C}_V}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \left(\frac{2V}{m_B+m_{K^*}} \epsilon_{\mu\nu\rho\sigma} \epsilon^\nu p^\rho p'^\sigma \right) \right\} O_{\mathcal{U}}^\mu. \tag{2.15}$$

And therefore the differential decay rate will be:

$$\frac{d\Gamma^{V\mathcal{U}}}{dE_{K^*}} = \frac{1}{8\pi^2 m_B} (q^2)^{d_{\mathcal{U}}-2} \sqrt{E_{K^*}^2 - m_{K^*}^2} \frac{A_{d_{\mathcal{U}}}}{\left(\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}\right)^2} \left\{ 8|C_V|^2 m_B^2 (E_{K^*}^2 - m_{K^*}^2) \frac{V^2}{(m_B+m_{K^*})^2} \right. \\ + |C_A|^2 \frac{1}{m_{K^*}^2 (m_B + m_{K^*})^2 q^2} \Big[(m_B+m_{K^*})^4 (3m_{K^*}^4 + 2m_B^2 m_{K^*}^2 - 6m_B m_{K^*}^2 E_{K^*} + m_B^2 E_{K^*}^2) A_1^2 \\ \left. + 4m_B^4 (E_{K^*}^2 - m_{K^*}^2)^2 A_2^2 + 4(m_B+m_{K^*})^2 (m_B E_{K^*} - m_{K^*}^2)(m_{K^*}^2 - E_{K^*}^2) m_B^2 A_1 A_2 \Big] \right\}. \tag{2.16}$$

To obtain the total decay width for $B \to K\mathcal{U}$ we must integrate over E_K in the range $m_K < E_K < (m_B^2 + m_K^2)/2m_B$, whereas to obtain the total decay width for $B \to K^*\mathcal{U}$ we must integrate over E_{K^*} in the range $m_{K^*} < E_{K^*} < (m_B^2 + m_{K^*}^2)/2m_B$.

3. Numerical results and conclusions

The total contribution to $B \to K(K^*) + \not{E}$ can be written as:

$$\Gamma = \Gamma^{\rm SM} + \Gamma^{\mathcal{U}} \ , \tag{3.1}$$

where the $\Gamma^{\rm SM}$ and $\Gamma^{\mathcal{U}}$ are the SM and unparticle contributions to the $B \to K(K^*) + \not{E}$ decay. And we should note that in the SM the missing energy in the final state is attributed to the presence of neutrinos. Hence the SM contribution to this process is given by $B \to K(K^*)\nu\bar{\nu}$. In the present case this signature can be mimicked by the process $B \to K(K^*)\,\mathcal{U}$, where we shall now try to estimate the bounds on the unparticles from the experimental constraints on missing energy signatures, as given by the B-factories BELLE and BaBar [17, 18]:

$$Br(B \to K\nu\bar{\nu}) < 1.4 \times 10^{-5} \ , \\ Br(B \to K^*\nu\bar{\nu}) < 1.4 \times 10^{-4} \ .$$

It is important to note that the SM process $B \to K(K^*)\nu\bar{\nu}$ provides a unique energy distribution spectrum of final state hadrons (K/K^* in our case). Presently the experimental limits on the branching ratio of these processes are about one order below the respective SM expectation values. However, these processes are expected to be measured at future SuperB factories [19]. As such, we presently only have an upper limit on the branching ratio of these processes, where to estimate the constraints on the unparticle properties.

Note that H. Georgi, in his first paper on unparticles, tried to emphasize that unparticles behave as a non-integral number of particles [6]. He further went on to analyze the

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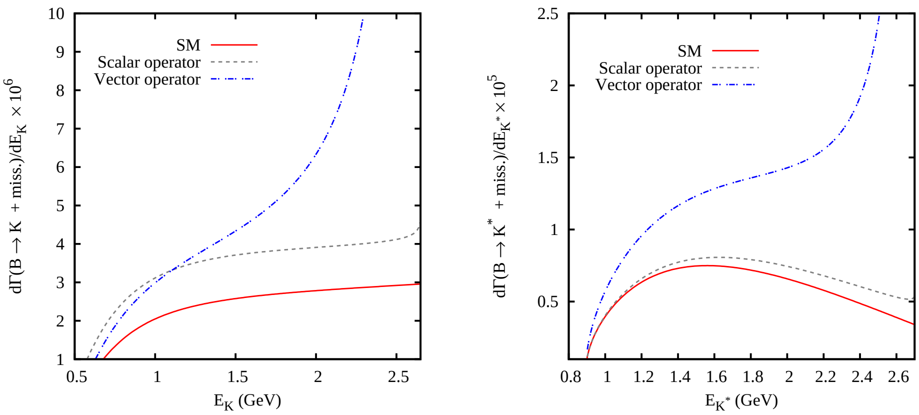

Figure 1: The differential branching ratio for: (a) Left panel: $B \to K + \not{E}$ as a function of the hadronic energy (E_K). (b) Right panel: $B \to K^* + \not{E}$ as a function of the hadronic energy (E_{K^*}). The other parameters are $d_{\mathcal{U}} = 1.9$, $\Lambda_{\mathcal{U}} = 1000$GeV, $\mathcal{C}_P = \mathcal{C}_S = 2 \times 10^{-3}$ and $\mathcal{C}_V = \mathcal{C}_A = 10^{-5}$.

distribution of the u-quark in the decay $t \to u\mathcal{U}$. It was argued that the peculiar shape of the distributions of E_u (the energy of the u-quark) may allow us to discover unparticles experimentally. As such, we have attempted to extend this same analogy to the process presently under consideration.

Finally, before presenting our numerical results, note that the future SuperB factories will be measuring the process $B \to K(K^*) + \not{E}$ by analyzing the spectra of the final state hadron. In doing this measurement at B-factories a cut for high momentum on the hadron is imposed, in order to suppress the background. Recall that unparticles would give us an unique distribution for the high energy hadron in the final state, such that in future B-factories one will be able to distinguish the presence of a scale invariant sector (or unparticles) by observing the spectrum of final state hadrons in $B \to K(K^*) + \not{E}$.

With this idea in mind we have tried to plot the differential decay width of $B \to K(K^*) + \not{E}$ as a function of $E_K(E_{K^*})$ in figure (1). As we can see from these figures the unparticle operators (especially the vector operators) give us a very distinctive distribution for the final state hadron's energy. The distribution of the unparticle contribution is strikingly different when we include a vector operator for a highly energetic final state hadron. As such, unparticle stuff can give a distinctly different signature from the SM in this regime, which it should be noted is experimentally more favorable at future SuperB factories.

In the next set of figures, figure (2), we have tried to analyze the constraints on the unparticle's scaling dimensions ($d_{\mathcal{U}}$) from different values of the cut-off scale $\Lambda_{\mathcal{U}}$. In these plots we have used some specific values of the effective couplings $\mathcal{C}_S$, $\mathcal{C}_P$, $\mathcal{C}_V$ and $\mathcal{C}_A$. As we can see from these figures the branching ratio is very sensitive to the scale dimension $d_{\mathcal{U}}$ and $\Lambda_{\mathcal{U}}$. In figure (3) we have shown the same plots for $B \to K^* + \not{E}$. From these two figures we can observe that the vector operators are more strongly constrained as compared to scalar operators. The second feature is that $B \to K + \not{E}$ provides better constraints

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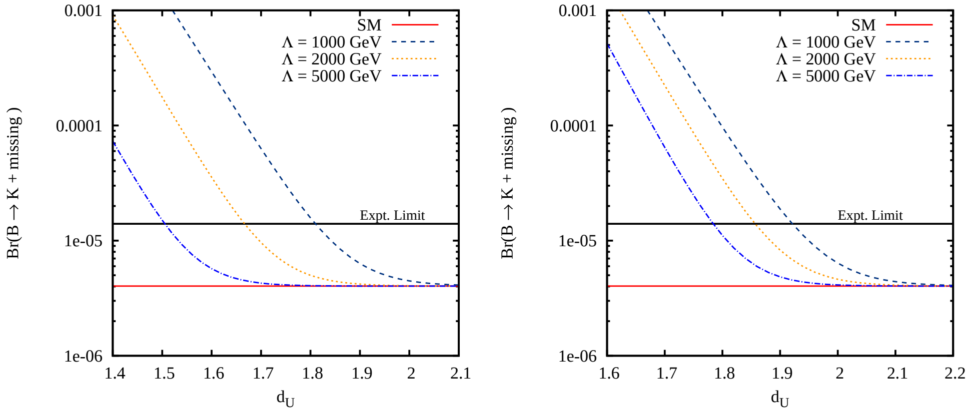

Figure 2: The branching ratio for $B \to K + \not{E}$ as a function of $d_{\mathcal{U}}$ for various values of $\Lambda_{\mathcal{U}}$. The left panel is for the contribution from the scalar operator, and the right panel is for the vector operator. The other parameters are $\mathcal{C}_S = 2 \times 10^{-3}$ and $\mathcal{C}_V = 10^{-5}$.

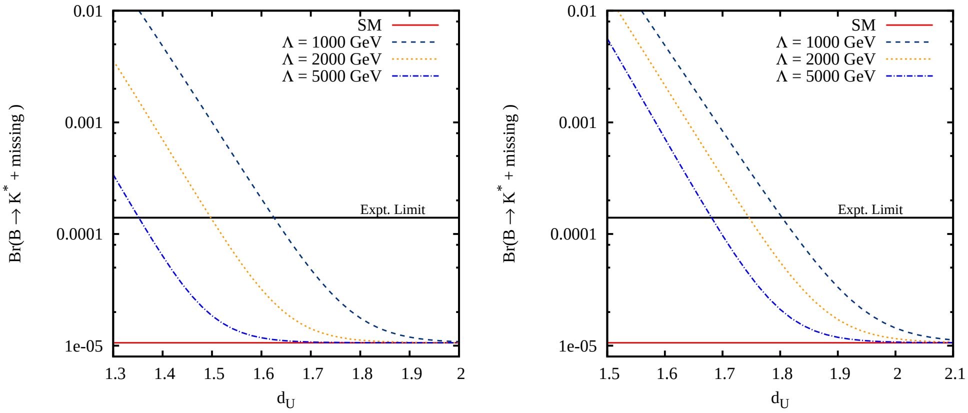

Figure 3: The branching ratio for $B \to K^* + \not{E}$ as a function of $d_{\mathcal{U}}$ for various values of $\Lambda_{\mathcal{U}}$. The left panel is for the contribution from the scalar operator, and the right panel is for the vector operator. The other parameters are $\mathcal{C}_P = 2 \times 10^{-3}$ and $\mathcal{C}_V = \mathcal{C}_A = 10^{-5}$.

than the $B \to K^* + \not{E}$ decay.

We have next tried to estimate the limits on the allowed values of the effective couplings, $\mathcal{C}_S$, $\mathcal{C}_P$, $\mathcal{C}_V$ and $\mathcal{C}_A$, from the present experimental limits on the branching ratio of $B \to K(K^*) + \not{E}$. Therefore, in figure (4) we have shown the dependence of the branching ratio of $B \to K + \not{E}$ on $\mathcal{C}_S$ and $\mathcal{C}_V$. As we can see from the expressions of the differential decay rate for $B \to K + \not{E}$, given in the previous section, if we consider the scalar (vector) operators, then the rate for this process is only dependent on $\mathcal{C}_S$ ($\mathcal{C}_V$).

Finally, in figure (5) we have shown the dependence of the branching ratio of $B \to K^* + \not{E}$ on the effective vertices. If we consider scalar operators then the rate of this process

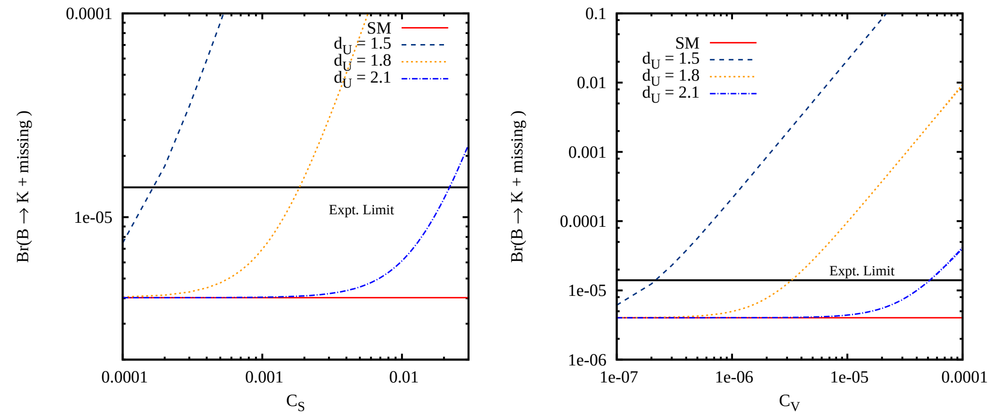

Figure 4: The branching ratio for $B \to K + \not{E}$ as a function of $\mathcal{C}_S$ (left panel) and $\mathcal{C}_V$ (right panel). The cutoff scale has been taken to be $\Lambda_{\mathcal{U}} = 1000\text{GeV}$.

is only dependent upon $\mathcal{C}_P$, whereas if we consider the vector operators then the rate can depend upon both $\mathcal{C}_V$ and $\mathcal{C}_A$.

To re-emphasize these last few points:

- $B \to K+$ scalar unparticle operator shall constrain the parameter $\mathcal{C}_S$,
- $B \to K^*+$ scalar unparticle operator shall constrain $\mathcal{C}_P$,
- $B \to K+$ vector unparticle operator will constrain only $\mathcal{C}_V$,
- whilst $B \to K^*+$ vector unparticle operator will constrain both $\mathcal{C}_V$ and $\mathcal{C}_A$.

To conclude, in this work we have analyzed the effects of unparticles on the missing energy signatures of rare B-decays. We have tried to argue that $B \to K(K^*) + \not{E}$ provides very useful constraints on the parameters of the model, where we have considered four operators, namely the scalar, pseudo-scalar, vector and axial vector operators. Both the modes $B \to K + \not{E}$ and $B \to K^* + \not{E}$ are different functions of these four operators, and hence provide independent constraints on the parameter space of the model. The set of operators we have considered, in principal, also contributes to meson anti-meson mixing, namely, $K - \bar{K}$, $B_d - \bar{B}_d$, $B_s - \bar{B}_s$ and $D - \bar{D}$ mixing. Detailed analyses of these, within the context of unparticle physics, has been done in reference [12]. Finally, note that the constraints provided by $B \to K(K^*) + \not{E}$ in some regions can be much stronger than the ones provided by meson anti-meson mixing.

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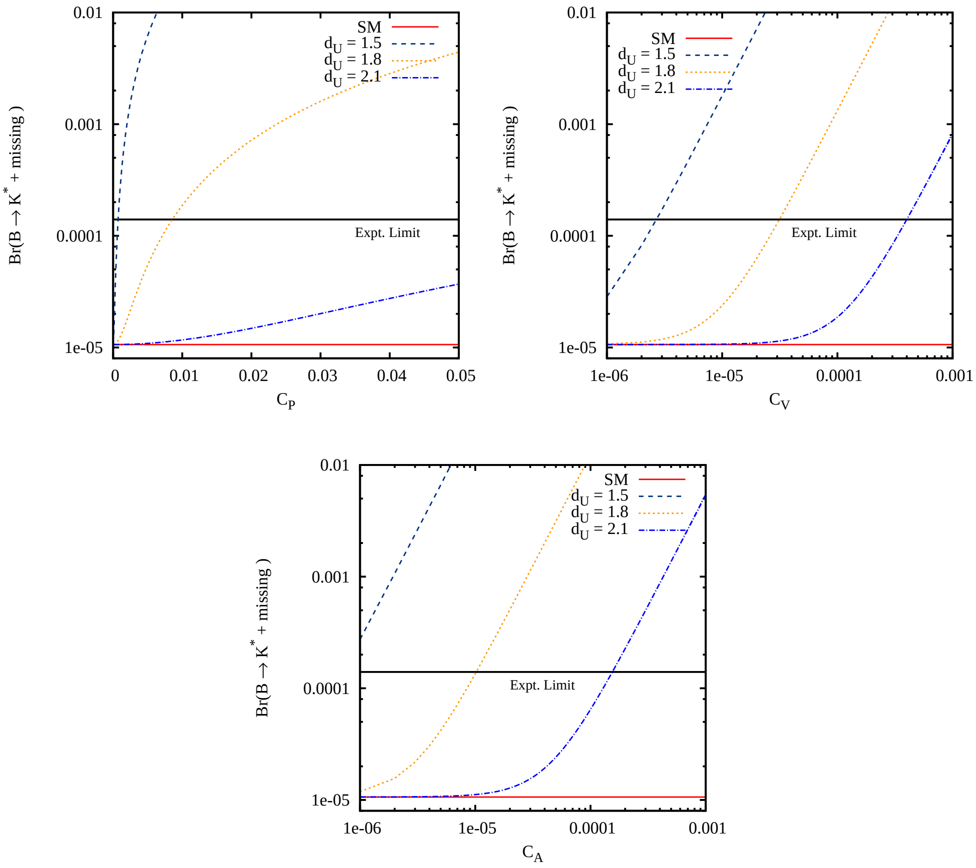

Figure 5: The branching ratio for $B \to K^* + \not{E}$ as a function of $\mathcal{C}_P$ (top left panel), $\mathcal{C}_V$ (top right panel) and $\mathcal{C}_A$ (bottom panel). The cutoff scale has been taken to be $\Lambda_{\mathcal{U}} = 1000\text{GeV}$.

A. The form factors

A.1 The form factors for the $B \to K$ transition

The form factors for the $B \to K$ transition can be written as [20]:

$$\begin{aligned} \langle K(p')|\bar{s}\gamma_\mu b|B(p)\rangle &= (p+p')_\mu f_+ + q_\mu f_- \ , \\ \langle K(p')|\bar{s}\gamma_\mu \gamma_5 b|B(p)\rangle &= 0 \ , \end{aligned} \tag{A.1}$$

where $q = p - p'$. Or alternately from the light cone sum rules [21] as:

$$\langle K(p')|\bar{s}\gamma_\mu b|B(p)\rangle = \left\{(p+p')_\mu - \frac{m_B^2 - m_K^2}{q^2} q_\mu\right\} f_+^P(q^2) + \left\{\frac{m_B^2 - m_K^2}{q^2} q_\mu\right\} f_-^P(q^2) \ . \tag{A.2}$$

Note that we can relate these two sets of form factors by:

$$\begin{aligned} f_+ &= f_+^P \ , \\ f_- &= \frac{m_B^2 - m_K^2}{q^2}\left(f_0^P - f_+^P\right) \ . \end{aligned} \tag{A.3}$$

	m_1	r_1	r_2	$m_{\rm fit}$
f_+^P	5.41	0.1616	0.1730	-
f_0^P	-	-	0.3302	5.41

Table 1: The parameters for the $B \to K$ form factors [20].

	r_1	r_2	$m_{\rm fit}^2$	m_R
V	0.923	- 0.511	49.40	5.32
A_0	1.364	- 0.99	36.78	5.63
A_1	-	0.290	40.38	-
A_2	-0.084	0.342	52.00	-

Table 2: The parameters for the $B \to K^*$ form factors [20].

In our numerical results we have followed the parameterization of Ball and Zwicky [21]:

$$\begin{aligned} f_0^P &= \frac{r_2}{1-q^2/m_{\rm fit}^2} \ , \\ f_+^P &= \frac{r_1}{1-q^2/m_1^2} + \frac{r_2}{(1-q^2/m_1^2)^2} \ , \end{aligned} \tag{A.4}$$

where the fitted parameters are given in table 1.

A.2 The form factors for the $B \to K^*$ transition

The form factors for the $B \to K^*$ transition can be written as [20]:

$$\begin{aligned} \langle K^*(p')|\bar{s}\gamma_\mu b|B(p)\rangle &= \epsilon_{\mu\nu\rho\sigma}\epsilon^\nu p^\rho p'^\sigma \frac{2V(q^2)}{m_B+m_{K^*}} \ , \\ \langle K^*(p')|\bar{s}\gamma_\mu\gamma_5 b|B(p)\rangle &= i\epsilon_\mu(m_B+m_{K^*})A_1(q^2) - i(p+p')_\mu(\epsilon.q)\frac{A_2(q^2)}{m_B+m_{K^*}} \\ &\quad -iq_\mu(\epsilon.q)\frac{2m_{K^*}}{q^2}[A_3(q^2)-A_0(q^2)] \ , \end{aligned} \tag{A.5}$$

where have again defined $q = p - p'$. For this transition we have used the parameterization of reference [20]:

$$\begin{aligned} F(q^2) &= \frac{r_1}{1-q^2/m_R^2} + \frac{r_2}{1-q^2/m_{\rm fit}^2} \ , \qquad \text{(for V, A}_0) \\ F(q^2) &= \frac{r_1}{1-q^2/m_{\rm fit}^2} + \frac{r_2}{(1-q^2/m_{\rm fit}^2)^2} \ , \qquad \text{(for A}_2) \\ F(q^2) &= \frac{r_2}{1-q^2/m_{\rm fit}^2} \ , \qquad \text{(for A}_1) \end{aligned} \tag{A.6}$$

where

$$A_3(q^2) = \frac{m_B+m_{K^*}}{2m_{K^*}}A_1(q^2) - \frac{m_B-m_{K^*}}{2m_{K^*}}A_2(q^2) \ . \tag{A.7}$$

Note that the fitted parameters used in the above equations have been given in table 2.

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